PHASE TRANSITIONS IN NEUTRON STARS

NORMAN K. GLENDENNING

Nuclear Science Division and
Institute for Nuclear & Particle Astrophysics
Lawrence Berkeley National Laboratory
Berkeley, California 94720
USA

Neutron stars are the densest objects in the universe today in which matter with several phases in adiabatic equilibrium can be found. Various high-density phases, both geometric and constitutional are spatially spread out by the pressure gradient in the star. Boundaries between phases slowly move, appear, or disappear as the density profile of the star is changed by the centrifugal force due to spindown caused by the magnetic torque of a pulsar, or the spinup of an x-ray neutron star because of the torque applied by mass accreted from a companion star. Phase transitions in turn produce their own imprint on the spin behavior through changes in the moment of inertia as one phase replaces another, in some cases on single stars, and in others on populations. These are the clues that we elucidate after first reviewing high-density phases.

1. A Brief History of Neutron Stars

- 1054 Chinese astronomer "observed the apparition of a guest star ...its color an iridescent yellow".
- 1933 Baade and Zwicky—binding energy of "closely packed neutrons" powers supernova.
- 1939 Oppenheimer, Volkoff and Tolman—neutron fermi gas.
- 1967 Pacini predicted magnetic dipole radiation.
- 1967 Hewish & Bell's serendipitous discovery of neutron stars producing a radio pulse once every revolution from beamed radiation along the magnetic axis which is fixed in the star. They are believed to be the *direct* product of core collapse a mature massive star and its the subsequent supernova.
- 1974 Hulse and Taylor binary neutron star pair in close orbit.
- 1984 Bacher's discovery of first Millisecond pulsar. They are believed to be very old supernova products that have been spun up by mass

accretion from a low-mass companion star.

- 1992 Wolszczan & Frail, discovery of 3 planets around a neutron star.

2. Gross Features of Neutron Stars

- Surface gravity M/R of Black hole =0.5,
  Neutron star =0.2,
  Sun $=10^{-6}$
- Gravitational binding / Nuclear binding ~ 10
- Radius $= 10 - 12$km, Mass$\geq 1.44 M_{\odot}$
- Spin periods from seconds to milliseconds
- Neutron stars are degenerate objects ($\mu << T$).
- Stars are electrically neutral. ($Z_{\rm Net}/A \sim (m/e)^2 < 10^{-36}$)
- Baryon number and charge are conserved.
- Strangeness not conserved (beyond 10^{-10} seconds).
- Millisecond pulsars have remarkably stable pulses:
  $P = 1.55780644887275 \pm 0.00000000000003$ ms
  (measured for PSR 1937+21 on 29 Nov 1982 at 1903 UT)

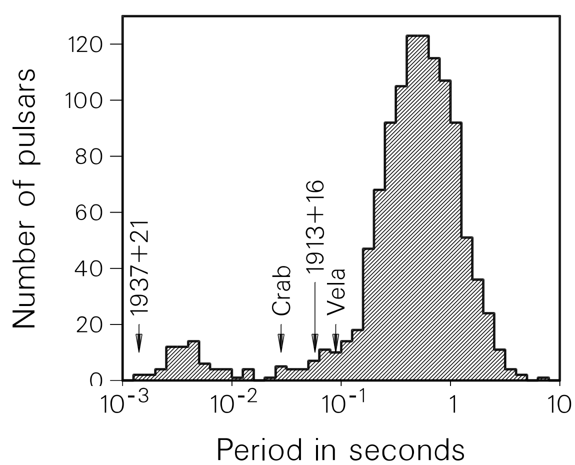


Figure 1. There are two classes of pulsars. The great bulk of known ones are the canonical pulsars with periods centered at about 0.7 seconds. The millisecond pulsars are believed to be an evolutionarily different class. They are harder to detect, and were first discovered in 1982.

3. Hyperonization

Free neutrons are unstable, but in a star the size and mass of a neutron star, *gravitational* binding energy is about ten times greater than *nuclear* binding so that neutrons are a stable component of dense stars. What about protons? The repulsive Coulomb force is so much stronger than the gravitational, that the net electric charge on a star must be very small ($Z_{\rm net}/(N+Z) < (m/e)^2 \sim 10^{-36}$). We can say that it is charge neutral. Since $m_p + m_e > m_n$, neutrons are the preferred baryon species. However, being Fermions, with increasing density of neutron matter, the Fermi level of neutrons will exceed the mass of proton and electron at some, not too high a density. Therefore, protons and electrons will also occupy neutron star matter. Because strangeness is conserved only on a weak interaction time-scale, this quantum number is not conserved in an equilibrium state. So with increasing density, the Pauli principle assures us that baryons of many species will be ingredients of dense neutral matter.[1,2]

Generally, it suffices to take the baryon octet into account together with electrons and muons. In Figure 2 we see that the Λ is most strongly populated in the center of a typical neutron star *if* quarks have not become deconfined at those densities. Notice that the lepton populations decrease as the populations of negatively charged hyperons increase. This is in accord with conservation of baryon number in the star. The number of electrons and muons are not by themselves conserved.

The equation of state is softened in comparison with a neutron matter equation of state. The softening means that the Fermi pressure is reduced so that hyperon matter cannot support as large a mass against gravitational collapse than would be the case otherwise. The hyperon transition is second order; particle populations vary continuously with density in a uniform medium. However, the densities reached in neutron star cores, 5 to 10 times nuclear matter density, are in all likelihood too high for baryons to exist as separate entities—quarks are likely to become deconfined at lower density than that. This is likely to be a first order phase transition.

4. First Order Transitions in Stars

Generally, physicists think of a phase transition such as from water to vapor as being *typical* of a first order transition. In the real world it is far from typical. Its characteristics are: if heated at constant pressure, the temperature of water and vapor will rise to 100 C and remain there until all the water has been evaporated before the temperature of the steam rises. This

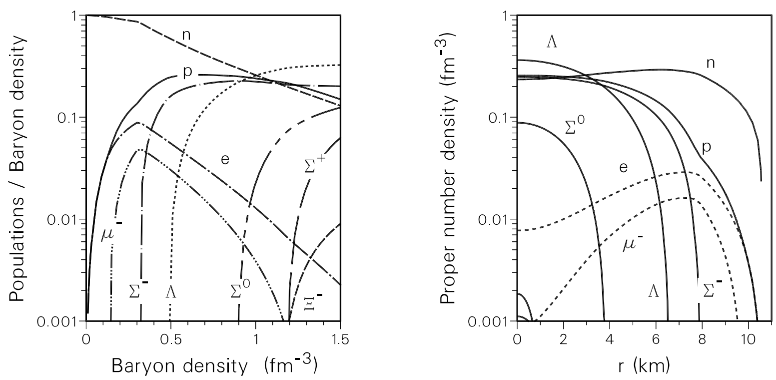


Figure 2. Particle populations as a function of baryon density in dense matter, and as a function of radial coordinate in a neutron star.

is true of substances having *one* independent component (like H^2O). The situation can, and usually is much more interesting for substances with two or more independent components, as I recognized a few years ago.[3] Neutron stars are an example. The independent components are the conserved baryon and electric charge. Until about 1990, all authors forced stellar models into the mold of single-component substances by imposing a condition of *local* charge neutrality and ignoring the discontinuity in electron chemical potential at the interface of two phases in equilibrium. In 1992 I realized that all these models of phase transitions in nuclear matter—which has two independent conserved components, the total baryon charge and the electric charge—were intrinsically incorrect.[3] They cannot satisfy Gibbs criteria for phase equilibrium in complex systems. And I stressed that an altogether new set of phenomenon were introduced by solving the problem correctly. Indeed, a *Coulomb crystalline region* involving the two phases in equilibrium could form, an idea that had not previously come to light.[3]

4.1. Degrees of freedom and driving forces

Two features can come into play in phase transition of complex substances that are absent in simple substances. The degree(s) of freedom can be seen in the following way. Imagine assembling a star in a pure phase (say ordinary nuclear matter) with B baryons and Q electric charges, either positive, negative or zero. (Of course, more precisely, we consider a typical local inertial region.) The concentration is said to be $c = Q/B$. Now consider another local region deeper in the star and at higher pressure with the same number of baryons and charges, but with conditions such that part of the volume is in the first phase and another part in the second phase. Suppose the baryons and charges in the two phases are distributed such that concentrations in the two phases are

$$Q_1/B_1 = c_1 \quad \text{and} \quad Q_2/B_2 = c_2 \,.$$

The conservation laws are still satisfied if

$$Q_1 + Q_2 = Q, \qquad B_1 + B_2 = B \,.$$

Why might the concentrations in the two phases be different from each other and from the concentration in the other local volumes at different pressure? Because the *degree of freedom* of redistributing the concentration may be exploited by *internal forces* of the substance so as to achieve a lower free energy. In a single-component substance there was no such degree of freedom, and in an n-component substance there are $n-1$ degrees of freedom. In deeper regions of the star, still different concentrations may be favored in the two phases in equilibrium at these higher-pressure locations. So you see that the each phase in equilibrium with the other, may have continuously changing properties from one region of the star to another. (This is unlike the simple substance whose properties remain unchanged in each equilibrium phase, until only one phase remains.)

The key recognition is that conserved quantities (or independent components) of a substance are conserved *globally*, but need not be conserved *locally*.[3] Otherwise, Gibbs conditions for phase equilibrium cannot be satisfied. Let us see how this is done.

Gibbs condition for phase equilibrium in the case of two conserved quantities is

$$p_1(\mu_n, \mu_e, T) = p_2(\mu_n, \mu_e, T)$$

We have introduced the neutron and electron chemical potentials by which baryon and electric charge conservation are to be enforced. In contrast to the case of a simple substance, for which Gibbs condition—$p_1(\mu, T) = p_2(\mu, T)$—can be solved for μ, the phase equilibrium condition cannot be satisfied for substance of more than one independent component without additional conservation constraints. Clearly, *local* charge conservation ($q(r) \equiv 0$) must be abandoned in favor of global conservation ($\int q(r)q(r) \equiv 0$), which is after all what is required by physics. For a uniform distribution global neutrality reads,

$$(1-\chi)q_1(\mu_n, \mu_e, T) + \chi q_2(\mu_n, \mu_e, T) = 0 \,,$$

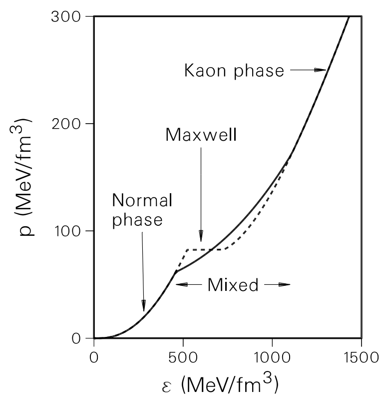


Figure 3. Solid line: equation of state for neutron star matter with a kaon condensed phase. Regions of the normal nuclear matter phase, the mixed phase, and the pure kaon condensed phase are marked. Notice that the pressure changes monotonically through the mixed phase. Dashed line: The Maxwell construction with the typical constant pressure region does not satisfy equality of the electron chemical potential in the two phases.

where $\chi = V_2/V,\ V = V_1 + V_2$. Given T and χ we can solve for $\mu_n and \mu_e$. Thus the solutions are of the form

$$\mu_n = \mu_n(\chi, T), \qquad \mu_e = \mu_e(\chi, T)\,.$$

Because of the dependance on χ, we learn that all properties of the phases in equilibrium change with proportion, χ, of the phases. This contrasts with simple (one component) substances. These properties are illustrated for the pressure in Figure 3. Behavior of the pressure is illustrated for two cases: (1) a simple, and (2) a complex substance. In the latter case, the pressure is monotonic, as proven above. This is in marked contrast to the pressure plateau of the simple (one component) substance.

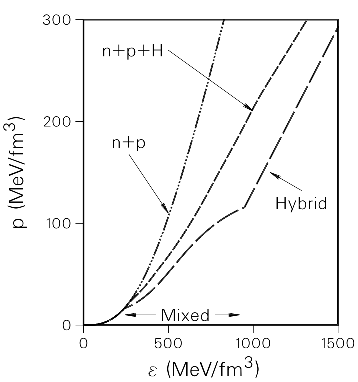
4.2. Isospin symmetry energy as a driving force

A well known feature of nuclear systematics is the valley of beta stability which, aside from the Coulomb repulsion, endows nuclei with $N = Z$ the greatest binding among isotones ($N + Z = \text{const}$). Empirically, the form of the symmetry energy is

$$E_{\rm N-sym} = -\epsilon[(N-Z)/(N+Z)]^2\,.$$

Physically, this arises in about equal parts from the difference in energies of neutron and proton Fermi energies and the coupling of the ρ meson to nucleon isospin current. Consider now a neutron star. While containing many nucleon species, neutron star matter is still very isospin asymmetric—it sits high up from the valley floor of beta stability—and must do so because of the asymmetry imposed by the strength of the Coulomb force compared to the gravitational.

Let us examine sample volumes of matter at ever-deeper depth in a star until we arrive at a local inertial volume where the pressure is high enough that some of the quarks have become deconfined; that both phases are present in the local volume. According to what has been said above, the highly unfavorable isospin of the nuclear phase can lower its repulsive asymmetry energy if some neutrons exchange one of their d quarks with a u quark in the quark phase in equilibrium with it. In this way the nuclear matter will become positively charged and the quark matter will carry a compensating negative charge, and the overall energy will be lowered. The *degree* to which the exchange will take place will vary according to the proportion of the phases—clearly a region with a small proportion of quark matter cannot as effectively relieve the isospin asymmetry of a large



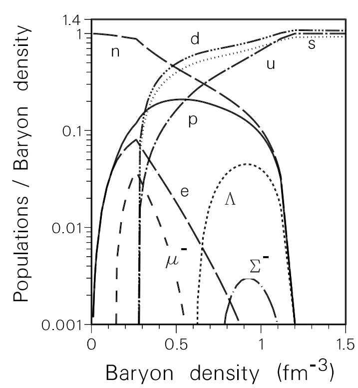


Figure 4. Equation of state for matter in *beta equilibrium* for three hypothetical models of dense nuclear matter; (1): only neutrons and protons are present $(\mathrm{n}+\mathrm{p})$, (2): in addition to neutrons and protons, hyperons (H) are also present $(\mathrm{n}+\mathrm{p}+\mathrm{H})$, (3): *Hybrid* denotes the equation of state for which matter has a low density nuclear phase, an intermediate mixed phase, and a high-density quark phase. Discontinuities in slope signal the transition between these phases.

Figure 5. The particle populations are shown as a function of density as phases change. The low-density region, $0.3\mathrm{fm}^{-3}$ ρ_B, is pure charge-neutral nuclear matter; the mixed nuclear and quark matter region lies in the density range $0.3 < \rho_B < 1.2$ fm^{-3}, and pure quark region lies above $\rho_{\mathrm{B}} > 1.2$ fm^{-3} .

proportion of neutron star matter of its excess isospin as can a volume of the star where the two phases are in more equal proportion. We see this quantitatively in Figure 6 where the charge densities on hadronic and quark matter are shown as a function of proportion of the phases.

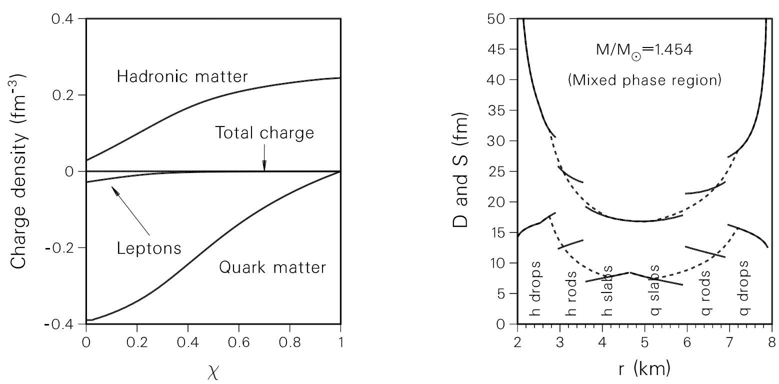


Figure 6. Charge densities on Hadronic and Quark matter as a function of proportion. Note that overall the mixture is neutral.

Figure 7. Diameter (bottom curve) and spacing (top curve) of the geometrical phases are shown as a function of position r in the star of $1.454M_{\odot}$. (see also Figure 8)

4.3. Geometrical phases

In equilibrium, the isospin driving force tends to concentrate positive charge on nuclear matter and compensating negative charge on quark matter. The Coulomb force will tend to break up regions of like charge while the surface interface energy will resist this tendency. The same competition is in play in the crust of the star where ionized atoms sit at lattice sites in an electron sea. For the idealized geometries of spheres, rods, or sheets of the rare phase immersed in the dominant one, and employing the Wigner-Seitz approximation (in which each cell has zero total charge, and does not interact with other cells), closed form solutions exist for the diameter D, and spacing S of the Coulomb lattice. The Coulomb and surface energy for drops, rods or slabs ($d = 3, 2, 1$) have the form:

$$\epsilon_C = C_d(\chi)D^2 \,, \quad \epsilon_S = S_d(\chi)/D \,,$$

where C_d abd S_d are simple algebraic functions of χ. The sum is minimized by $\epsilon_S = 2\epsilon_C$. Hence, the diameter of the objects at the lattice sites is

$$D = [S_d(\chi)/2C_d(\chi)]^{1/3} \, ,$$

where their spacing is $S = D/\chi^{1/d}$ if the hadronic phase is the background or $S = D/(1-\chi)^{1/d}$ if the quark phase is background. Figure 7 shows the computed diameter and spacing of the various geometric phases of quark and hadronic matter as a function of radial coordinate in a hybrid neutron star.

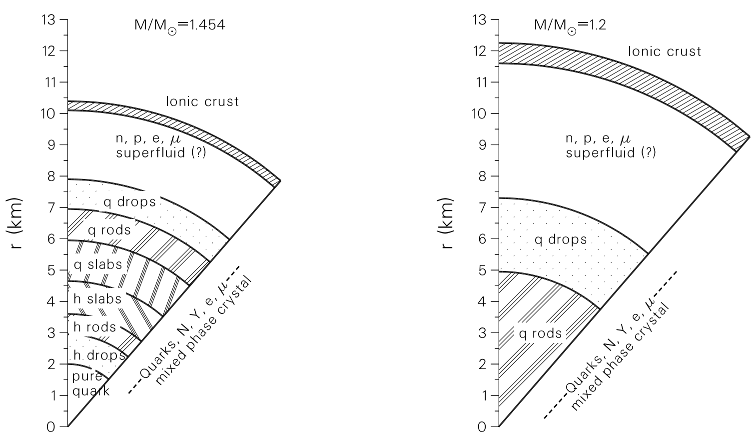


Figure 8. Pie sections showing geometric phases in two stars of different mass

4.4. Color-flavor locked quark-matter phase (CFL)

Rajagopal and Wilczek have argued that the Fermi surface of the quark deconfined phase is unstable to correlations of quarks of *opposite* momentum and *unlike* flavor and form BCS pairs[4]. They estimate a pairing gap of $\Delta \sim 100$ MeV. The greatest energy benefit is achieved if the Fermi surfaces of all flavors are equal in radius. This links color and flavor by an invariance to simultaneous rotations of color and flavor. The approximate energy density corresponding to the gap is

$$\epsilon_{\Delta-\mathrm{CFL}} \sim -C(k_F \, \Delta)^2 \sim 50 \cdot C \ \mathrm{MeV/fm}^3 \, ,$$

where C is an unknown constant. This is another "driving force" as spoken of above *in addition* to the nuclear symmetry energy ϵ_{sym}. It acts, not to restore isospin symmetry in nuclear matter, but color-flavor symmetry in the quark phase. Alford, Rajagopal, Reddy, and Wilczek have argued that the CFL phase, which is *identically* charge neutral and has this large pairing gap *may* preempt the possibility of phase equilibrium between confined hadronic matter and the quark phase; that any amount of quark matter would go into the charge neutral CFL phase (with equal numbers of u, d and s quarks, irrespective of mass) and that the mixed phase spoken of above would be absent.[5] That the nuclear symmetry driving force would be overcome by the color-flavor locking of the quark phase leaving the degree of freedom possessed by the two-component system unexploited. The discontinuity of the electron chemical potential in the two phases, hadronic and quark matter would be patched by a spherical interface separating a core of CFL phase in the star from the surrounding hadronic phase. For that conclusion to be true, a rather large surface interface coefficient was chosen by dimensional arguments.

However, my opinion is that nature will make a choice of surface interface properties between hadronic and quark matter such that the degree of freedom of exchanging charge can be exploited by the driving forces (here two in number as discussed below). This is usually the case. Physical systems generally have their free energy lowered when a degree of freedom (as spoken of above) becomes available.

With two possible phases of quark matter, the uniform uncorrelated one discussed first, and the CFL phase as discussed by Rajagopal and Wilczek, there is now a competition between the CFL pairing and the nuclear symmetry-energy densities, and these energy densities are *weighted* by the volume proportion χ of quark matter in comparison with hadronic matter in locally inertial regions of the star. That is to say, $\epsilon_{\rm CFL}$ and $\epsilon_{\rm sym}$ are not directly in competition, but rather they are weighted by the relevant volume proportions. It is not a question of "either, or" but "one, then the other".

The magnitude of the nuclear symmetry energy density at a typical phase transition density of $\rho \sim 1/{\rm fm}^3$ is

$$\epsilon_{\rm N-sym} = -35[(N-Z)/(N+Z)]^2 \ {\rm MeV/fm}^3 \, .$$

To gain this energy a certain price is exacted from the disturbance of the symmetry of the uniform quark matter phase in equilibrium with it; $\epsilon_{\rm Q-sym}$. As can be inferred from Figure 6, the price is small compared to the gain. On the other side, the energy gained by the quark matter entering the CFL phase was written above and is offset by the energy not gained by the nuclear matter because the CFL preempts an improvement in its isospin asymmetry. So we need to compare

$$(1-\chi)\epsilon_{\rm N-sym} - \chi\,\epsilon_{\rm Q-sym} - [\epsilon_{\rm surf}(\chi) + \epsilon_{\rm coul}(\chi)]$$

with

$$\chi\,\epsilon_{\Delta-\rm CFL} - (1-\chi)\epsilon_{\rm N-sym}\,.$$

The behavior of these two lines as a function of proportion of quark phase χ in a local volume in the star is as follows:[a] The first expression for the net gain in energy due to the formation of a mixed phase of nuclear and uniform quark matter monotonically decreases from its maximum value at $\chi = 0$ while the second expression, the net energy gain in forming the CFL phase monotonically increases from *zero* at $\chi = 0$. Therefore as a function of χ or equivalently depth in the star measured from the depth at which the first quarks become deconfined, nuclear symmetry energy is the dominating driving force, while at some value of χ in the range $0 < \chi < 1$ the CFL pairing becomes the dominating driving force.

In terms of Figure 8, several of the outermost geometric phases in which quark matter occupies lattice sites in a background of nuclear matter are undisturbed. But the sequence of geometric phases is terminated before the series is complete, and the inner core is entirely in the CFL phase.

In summary, when the interior density of a neutron star is sufficiently high as to deconfine quarks, a charge neutral color-flavor locked phase with no electrons will form the inner core. This will be surrounded by one or more shells of mixed phase of quark matter in a uniform phase in phase equilibrium with confined hadronic matter, the two arranged in a Coulomb lattice which differs in dimensionality from one shell to another. As seen in Figure 6, the density of electrons is very low to essentially vanishing, because overall charge neutrality can be achieved more economically among the conserved baryon charge carrying particles. Finally, All this will be surrounded by uniform charge neutral nuclear matter with varying particle composition according to depth (pressure), (cf. Figure 5.)

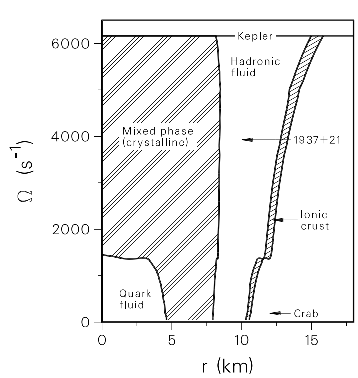
5. Rotation and Phase Transitions

Except for the first few seconds in the life of a neutron star, at which time they radiate the vast bulk of their binding energy in the form of neutrinos, we think of them as rather static objects. However the spin evolution at *millisecond* periods of rotation brings about centrifugally induced changes in the density profile of the star, and hence also in the thresholds and

[a]The behavior of the quantity in square brackets can be viewed in Figure 9.14 of reference [[2], 2'nd ed.]

densities of various hyperons, dense phases such as kaon condensed phase and inevitably quark matter. We shall *assume* that the central density of the more massive millisecond pulsars—being centrifugally diluted—lies below the critical density for pure quark matter, while the central density of canonical pulsars, like the Crab, and more slowly rotating ones, lies above. We explore the consequences of such assumptions.

Because of the different compressibility of low and high-density phases, conversion from one phase to another as the phase boundary slowly moves with changing stellar spin (Fig. 9) results in a considerable redistribution of mass (Fig. 10) and hence change in moment of inertia over time. The time scale is of the order of 10^7 to 10^9 yr. The behavior of the moment of inertia while successive shells in the star are changing phase is analogous to the so-called backbending behavior of the moment of inertia of deformed rotating nuclei brought about by a change of phase from one in which the coriolis force breaks nucleon spin pairing to one in which spins are paired. Compare Figs. 11 and 12.



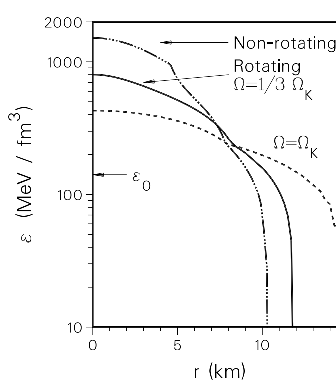


Figure 9. Radial boundaries at various rotational frequencies separating various phases. The frequencies of two pulsars, the Crab and PSR 1937+21 are marked for reference.

Figure 10. Mass profiles as a function of equatorial radius of a star rotating at three different frequencies. At low frequency the star is very dense in its core, having a 4 km central region of highly compressible pure quark matter. Inflections at $\epsilon \approx 220$ and 950 are the boundaries of the mixed phase.

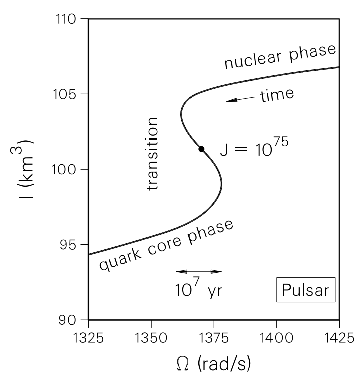
Elsewhere we have discussed the possible effect of a phase transition on isolated millisecond radio pulsars.[6] Here we discuss x-ray neutron stars that have a low-mass non-degenerate companion. Beginning at a late stage in the evolution of the companion it evolves toward its red-giant stage and mass overflows the gravitational barrier between the donor and neutron star. The neutron star is spun up by angular momentum conservation of the accreted matter. The heated surface of the neutron star and its rotation may be detected by emitted x-rays.

In either case—neutron star accretors or millisecond pulsars—the radial thresholds of particle types and phase boundaries will move—either outward or inward—depending on whether the star is being spun up or down. The critical density separating phases moves slowly so that the conversion from one phase to another occurs little by little at the moving boundary. In a rapidly rotating pulsar that is spinning down, the matter density initially is centrifugally diluted, but the density rises above the critical phase transition density as the star spins down. Relatively stiff nuclear matter is converted to highly compressible quark matter. The overlaying layer of nuclear matter squeezes the quark matter causing the interior density to rise, while the greater concentration of mass at the center acts further to concentrate the mass of the star. Therefore, its moment of inertia decreases over and above what would occur in an immutable rotating gravitating fluid that is spinning down. If this occurs, the moment of inertia as a function of spin exhibits a backbend as in Figure 11. Such a phenomenon has been observed in nuclei, as illustreated in Figure 12.[18,19,20] The opposite evolution of the moment of inertia may occur in x-ray neutron stars that are spinning up when the spin change spans the critical region of phase transition.

As a result of the backbend in moment of inertia, an isolated ms pulsar may cease its spindown and actually spin up for a time, even though loosing angular momentum to radiation as was discussed in a previous work.[6] An x-ray neutron star with a companion may pause in its accretion driven spinup until quark matter is driven out of the star, after which it will resume spinup. Spinup or spin down occurs very slowly, being controlled by the mass accretion rate or the magnitude of the magnetic dipole field, respectively. So, the spin anomaly that might be produced by a conversion of matter from one phase to another will endure for many millions of years. If it were fleeting it would be unobservable. But enduring for a long epoch—if the phenomenon occurs at all—it has a good chance of being observed.

A very interesting work by Spyrou and Stergioulas has recently appeared in the above connection.[7] They perform a more accurate numerical calculation for a rotating relativistic star, as compared to our perturbative solution. They find that the backbend in our particular example occurs very close to, or at the maximum (non-rotating) star, but that it is generic for stars that are conditionally stabilized by their spin. This is possibly the situation for some or eventually all accretors.

In fact, we expect a phase transition to leave a *permanent* imprint on the distribution in spins of x-ray accretors. Because of the increase of moment of inertia during the epoch in which the quark core is driven out of a neutron star as it is spun up by mass accretion, spinup is—during this epoch—hindered. Therefore we expect the population of accretors to be clustered in the spin-range corresponding to the expulsion of the quark phase from the stellar core. Spin clustering was reported in the population of x-ray neutron stars in binaries that were reported in data from the Rossi X-ray Timing Explorer.[8] However, later observations failed to confirm this.



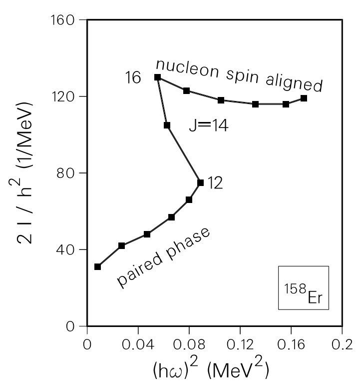


Figure 11. Development of moment of inertia of a model neutron star as a function of angular velocity. The backbend in this case is similar to what is observed in some rotating nuclei. (Adapted from Ref. [6].)

Figure 12. Backbending in the rotating Er nucleus and an number of others was discovered in the 1970s.

6. Calculation

The theory and parameters used to describe our model neutron star are precisely those used in previous publications. Its initial mass is $M = 1.42 M_{\odot}$, close to the mass limit. The confined hadronic phase is described by a generalization of a relativistic nuclear field theory solved at the mean field level in which members of the baryon octet are coupled to scalar, vector and vector-isovector mesons.[9,2] The parameters [6,10] of the nuclear Lagrangian were chosen so that symmetric nuclear matter has the following properties: binding energy $B/A = -16.3$ MeV, saturation density $\rho = 0.153\,\mathrm{fm}^{-3}$, compression modulus $K = 300$ MeV, symmetry energy coefficient $a_{\rm sym} = 32.5$ MeV, nucleon effective mass at saturation $m^{\star}_{\rm sat} = 0.7m$. These together with the ratio of hyperon to nucleon couplings of the three mesons, $x_\sigma = 0.6,\ x_\omega = 0.653 = x_\rho$ yield the correct Λ binding in nuclear matter.[10]

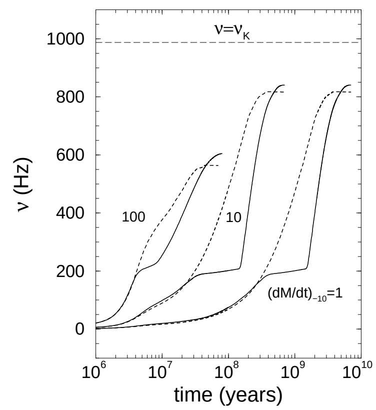
Quark matter is treated in a version of the MIT bag model with the three light flavor quarks ($m_u = m_d = 0,\ m_s = 150$ MeV) as described. [11] A value of the bag constant $B^{1/4} = 180$ MeV is employed.[6] The transition between these two phases of a medium with two independent conserved charges (baryon and electric) has been described elsewhere.[3] We use a simple schematic model of accretion. [12,13,14] All details of our calculation can be found elsewhere.[15,16,17]

7. Results

The spin evolution of accreting neutron stars as determined by the changing moment of inertia and the evolution equation[15] is shown in Fig. 13. We assume that up to $0.4M_\odot$ is accreted. Otherwise the maximum frequency attained is less than shown. Three *average* accretion rates are assumed, $\dot{M}_{-10} = 1$, 10 and 100 (where $\dot{M}_{-10}$ is in units of $10^{-10}M_\odot/\mathrm{y}$).

We compute a frequency distribution of x-ray stars in low-mass binaries (LMXBs) from Fig. 13, for *one* accretion rate, by assuming that neutron stars begin their accretion evolution at the average rate of one per million years. A different rate will only shift some neutron stars from one bin to an adjacent one. The donor masses in the binaries are believed to range between 0.1 and $0.4M_\odot$ and we assume a uniform distribution in this range and repeat the calculation shown in Fig. 13 at intervals of $0.1M_\odot$. The resulting frequency distribution of x-ray neutron stars is shown in Fig. 14; it is striking. A spike in the distribution signals spinout of the quark matter core as the neutron star spins up. This feature would be absent if there were no phase transition in our model of the neutron star.

The data in Fig. 14 is gathered from Tables 2–4 of the review article of van der Klis concerning discoveries made with the Rossi X-ray Timing Explorer.[8] However, later observations have failed to confirm the original report.



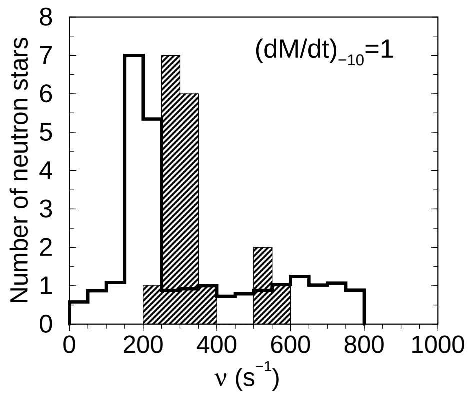


Figure 13. Evolution of spin frequencies of accreting neutron stars with (solid curves) and without (dashed curves) quark deconfinement if $0.4M_\odot$ is accreted. The spin plateau around 200 Hz signals the ongoing process of quark confinement in the stellar centers. Spin equilibrium is eventually reached. (From Ref. [15].)

Figure 14. Calculated spin distribution of the underlying population of x-ray neutron stars for one accretion rate (open histogram) is normalized to the number of observed objects (18) at the peak. Data on neutron stars in low-mass X-ray binaries (shaded histogram) is from Ref. [8]. These results have not been observed in later observations, however. The spike in the calculated distribution corresponds to the spinout of the quark matter phase. Otherwise the spike would be absent. (From Ref. [15].)

8. Conclusion

We find that if a clustering in rotation frequency of accreting x-ray neutron stars in low-mass binaries were discovered, it could be caused by the progressive conversion of quark matter in the core to confined hadronic matter, paced by the slow spinup due to mass accretion. When conversion is completed, normal accretion driven spinup resumes. To distinguish this conjecture from others, one would have to discover the inverse phenomenon—a spin anomaly near the same frequency in an isolated ms pulsar.[6] If such a discovery were made, and the apparent clustering of x-ray accretors is confirmed, we would have some degree of confidence in the hypothesis that a dense matter phase, most plausibly quark matter, exists from birth in the cores of canonical neutron stars, is spun out if the star has a companion from which it accretes matter, and later, having consumed its companion, resumes life as a millisecond radio pulsar and spins down.

Acknowledgments

This work was supported by the Director, Office of Energy Research, Office of High Energy and Nuclear Physics, Division of Nuclear Physics, of the U.S. Department of Energy under Contract DE-AC03-76SF00098.

References

1. N. K. Glendenning, Phys. Lett. **114B** (1982) 392; Astrophys. J. **293** (1985) 470.
2. N. K. Glendenning, *Compact Stars* (Springer–Verlag New York, 1'st ed. 1996, 2'nd ed. 2000).
3. N. K. Glendenning, Phys. Rev. D **46** (1992) 1274.
4. K. Rajagopal and F. Wilczek, Phys.Rev.Lett. **86** (2001) 3492.
5. M. Alford, K. Rajagopal, S. Reddy, and F. Wilczek, Phys.Rev. D **64** (2001) 074017.
6. N. K. Glendenning, S. Pei and F. Weber, Phys. Rev. Lett. **79** (1997) 1603.
7. N.K. Spyrou and N. Sterigoulas, Astron. & Astrophys., **395** (2002) 151.
8. M. van der Klis, Ann. Rev. Astron. Astrophys, **38** 717 (2000).
9. N. K. Glendenning, Astrophys. J. **293** (1985) 470.
10. N. K. Glendenning and S. A. Moszkowski, Phys. Rev. Lett. **67** (1991) 2414.
11. E. Farhi and R. L. Jaffe, Phys. Rev. D **30** (1984) 2379.
12. R. F. Elsner and F. K. Lamb, Astrophys. J. **215** (1977) 897.
13. P. Ghosh, F. K. Lamb and C. J. Pethick, Astrophys. J. **217** (1977) 578.
14. V. M. Lupinov, *Astrophysics of Neutron Stars*, (Springer-Verlag, New York, 1992.
15. N. K. Glendenning and F. Weber, Astrophys. J. Lett **559** (2001) L119.
16. N. K. Glendenning and F. Weber, astro-ph/0010336 (2000).
17. N. K. Glendenning and F. Weber, *Signal of quark deconfinement in millisecond pulsars and reconfinement in accreting x-ray neutron stars*, in Physics of Neutron Star Interiors, Ed. by Blaschke, Glendenning and Sedrakian (Springer-Verlag, Lecture Notes Series, 2001).
18. B. R. Mottelson and J. G. Valatin, Phys. Rev. Lett. **5** (1960) 511.
19. A. Johnson, H. Ryde and S. A. Hjorth, Nucl. Phys. **A179** (1972) 753.
20. F. S. Stephens and R. S. Simon, Nucl. Phys. **A183** (1972) 257.